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Phenomenological implications of light stop and higgsinos [1]

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Abstract

We examine the phenomenological implications of light $\tilde{t}_R$ and higgsinos in the Minimal Supersymmetric Standard Model, assuming $\tan^2\beta < m_t/m_b$ and heavy $\tilde{t}_L$ and gauginos. In this simplified setting, we study the contributions to Δm_{B_d}, ϵ_K, $BR(b \to s\gamma)$, $R_b \equiv \Gamma(Z \to b\bar{b})/\Gamma(Z \to \text{hadrons})$, $BR(t \to bW)$, and their interplay.

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1. If low-energy supersymmetry (for a review and references, see e.g. [1]) plays a role in the resolution of the naturalness problem of the Standard Model (SM), then the Minimal Supersymmetric Standard Model (MSSM) is the most plausible effective theory at the electroweak scale, and we should be close to the discovery of Higgs bosons and supersymmetric particles. The scalar partners of the top quark (two complex spin-0 fields, one for each chirality of the corresponding quark) and the fermionic partners of the gauge and Higgs bosons (two charged Dirac particles, or charginos, and four neutral Majorana particles, or neutralinos) are among the most likely candidates for an early discovery. A particularly attractive possibility is the existence of light $\tilde{t}_R$ and higgsinos, within the discovery reach of LEP2, as suggested by some MSSM fits to precision electroweak data (see [2] for the different points of view). The aim of the present paper is to elucidate the phenomenological implications of such a possibility, developing some of the observations already present in [2] in a more systematic and transparent way.

In the rest of this section, we present simplified expressions for the light sparticle masses in the limit of interest. In section 2, we introduce simplified MSSM analytical formulae for a number of physical observables, such as Δm_{B_d}, ϵ_K, $BR(b \to s\,\gamma)$, $R_b \equiv \Gamma(Z \to b\bar{b})/\Gamma(Z \to \text{hadrons})$, $BR(t \to bW)$. For each class of processes, we discuss how the existing experimental data constrain the MSSM parameter space. In section 3, we present our conclusions.

Assuming that the squark mass matrices can be diagonalized (in generation space) simultaneously with those of the corresponding quarks, the spectrum of the stop sector is described, in the usual MSSM notation and in the $(\tilde{t}_L, \tilde{t}_R)$ basis, by the 2×2 matrix

$$\mathcal{M}^2_{\tilde{t}} = \begin{pmatrix} m^2_{LL} & m^2_{LR} \\ m^2_{LR} & m^2_{RR} \end{pmatrix}$$
$$= \begin{pmatrix} m^2_{Q_3} + m_t^2 + \frac{1}{6}(4m_W^2 - m_Z^2)\cos 2\beta & m_t(A_t - \mu\cot\beta) \\ m_t(A_t - \mu\cot\beta) & m^2_{U_3} + m_t^2 + \frac{2}{3}(-m_W^2 + m_Z^2)\cos 2\beta \end{pmatrix} . \tag{1}$$

In the limit $m^2_{RR}, m^2_{LR} \ll m^2_{LL}$, the lightest stop eigenstate is predominantly a $\tilde{t}_R$, $\tilde{t}_2 = -\sin\theta_t \tilde{t}_L + \cos\theta_t \tilde{t}_R$, with $\theta_t \simeq m^2_{LR}/m^2_{LL} \ll 1$, and the corresponding mass eigenvalue is given by $m^2_{\tilde{t}_2} \simeq m^2_{RR} - (m^2_{LR})^2/m^2_{LL}$. The above situation could arise for example when $\tan\beta < m_t/m_b$, or $h_t \gg h_b$, since in that case the structure of the renormalization group equations favours $m_{U_3} < m_{Q_3} < m_{\tilde{q}}$, where $m_{\tilde{q}}$ is some average squark mass. It is also known [2] that it is easier to reconcile a light $\tilde{t}_R$ than a light $\tilde{t}_L$ with the stringent limits on the effective ρ parameter coming from the electroweak precision data .

The mass matrices in the chargino and neutralino sector read

$$\mathcal{M}_C = \begin{pmatrix} M_2 & \sqrt{2} m_W \sin\beta \\ \sqrt{2} m_W \cos\beta & \mu \end{pmatrix} , \tag{2}$$

and

$$\mathcal{M}_N = \begin{pmatrix} M_1 & 0 & -m_Z c_\beta s_W & m_Z s_\beta s_W \\ 0 & M_2 & m_Z c_\beta c_W & -m_Z s_\beta c_W \\ -m_Z c_\beta s_W & m_Z c_\beta c_W & 0 & -\mu \\ m_Z s_\beta s_W & -m_Z s_\beta c_W & -\mu & 0 \end{pmatrix} , \tag{3}$$

where $\mathcal{M}_N$ has been written in the $(-i\tilde{B}, -i\tilde{W}_3, \tilde{H}_1^0, \tilde{H}_2^0)$ basis, and $s_\beta \equiv \sin\beta$, $c_\beta \equiv \cos\beta$, $s_W \equiv \sin\theta_W$, $c_W \equiv \cos\theta_W$. An approximate Peccei-Quinn symmetry, recovered in the limit $\mu \to 0$, may originate the hierarchy $\mu \ll M_1, M_2$, which leads to one charged and two neutral higgsinos much lighter than the other mass eigenstates. In first approximation, we find the following three degenerate eigenstates: $\tilde{H}_S \equiv (\tilde{H}_1^0 + \tilde{H}_2^0)/\sqrt{2}$, $\tilde{H}_A \equiv (\tilde{H}_1^0 - \tilde{H}_2^0)/\sqrt{2}$, $\tilde{H}^\pm$, with eigenvalues $|m_{\tilde{H}_S}| = |m_{\tilde{H}_A}| = |m_{\tilde{H}^\pm}| = |\mu|$: this will be sufficient for most of the following considerations. For the discussion of chargino and neutralino decays, it is useful to go beyond this approximation, to see how the degeneracy is lifted. Assuming as usual $(M_2/M_1) \simeq (3/5)\cot^2\theta_W$, corresponding to universal gaugino masses at some grand-unification scale, and expanding in $1/M_2$, we find[1]

$$\left|m_{\tilde{H}_S}\right| = |\mu + \Delta_S| \, , \quad \left|m_{\tilde{H}_A}\right| = |\mu + \Delta_A| \, , \quad |m_{\tilde{H}^\pm}| = |\mu + \Delta_C| \, , \tag{4}$$

where

$$\Delta_S = (1 - \sin 2\beta)\frac{4}{5}\frac{m_W^2}{M_2} \, , \qquad \Delta_A = -(1 + \sin 2\beta)\frac{4}{5}\frac{m_W^2}{M_2} \, , \qquad \Delta_C = -\frac{m_W^2 \sin 2\beta}{M_2} \, , \tag{5}$$

in agreement with [3]. We then find the mass hierarchies

$$\begin{aligned} &|m_{\tilde{H}_A}| < |m_{\tilde{H}^\pm}| < |m_{\tilde{H}_S}| \quad (\mu M_2 > 0) \, , \\ &|m_{\tilde{H}_S}| < |m_{\tilde{H}^\pm}| < |m_{\tilde{H}_A}| \quad (\mu M_2 < 0) \, , \end{aligned} \tag{6}$$

consistent with the phenomenological request of a neutral and weakly interacting lightest supersymmetric particle. The typical size of the mass splittings, according to eq. (5), is illustrated in fig. 1. Since we are not assuming a large mixing in the stop sector, we expect radiative corrections to the previous formulae to be negligible [4].

2. In this section we present simplified analytical formulae describing the MSSM contributions to a number of important observables, in the special case of light $\tilde{t}_R$ and higgsinos, and we discuss the resulting phenomenological constraints on the associated parameter space. Before proceeding, we would like to state clearly the assumptions under which the following discussion will be valid: 1) $\tilde{t}_R$ and higgsinos are approximate mass eigenstates, with $\tilde{t}_L$ and gauginos sufficiently heavy to give negligible contributions; 2) $\tan^2\beta < m_t/m_b$,

[1] With a slight abuse of notation, we keep the symbols $\tilde{H}_S$, $\tilde{H}_A$ and $\tilde{H}^\pm$ also for the perturbed eigenstates.

which allows us to neglect the vertices proportional to the bottom Yukawa coupling h_b, with respect to those proportional to the top Yukawa coupling h_t ($\tan\beta < m_t/m_b$ would be sufficient for the stop and chargino couplings, whereas $\tan^2\beta < m_t/m_b$ will be required by our approximations for the charged Higgs couplings); 3) negligible flavour-changing effects associated with the quark-squark-gluino and the quark-squark-neutralino vertices. Since the theoretical expressions for the observables to be discussed below have a strong dependence on the top quark mass, we would like to recall here the one-loop QCD relation[2] between m_t, the $\overline{MS}$ running mass at the top-mass scale, and the pole mass M_t: $M_t = m_t[1 + (4/3)\alpha_s/\pi]$. For definiteness, we shall present our results for the input value $m_t = 170$ GeV (corresponding to $M_t \simeq 178$ GeV for $\alpha_s \simeq 0.12$), compatible with the present Tevatron data [6].

$\boxed{\Delta m_B,\, \epsilon_K}$

We discuss here the MSSM contributions to the B_d^0–$\overline{B_d^0}$ mass difference Δm_{B_d} and to the CP-violation parameter of the K^0-$\overline{K^0}$ system ϵ_K, and the constraints on the model parameters coming from the experimentally measured values of Δm_{B_d} and ϵ_K.

For our purposes, a convenient way of parametrizing the B_d^0–$\overline{B_d^0}$ mass difference is [7]:

$$\Delta m_{B_d} = \eta_{B_d} \cdot \frac{4}{3} f_{B_d}^2 B_{B_d} \cdot m_{B_d} \cdot \left(\frac{\alpha_W}{4 m_W}\right)^2 \cdot |K_{tb} K_{td}^*|^2 \cdot x_{tW} \cdot |\Delta| \,, \tag{7}$$

where $\eta_{B_d} \simeq 0.55$ is a QCD correction factor; f_{B_d} is the B_d decay constant and B_{B_d} the vacuum saturation parameter; $\alpha_W = g^2/(4\pi)$, K is the Kobayashi-Maskawa matrix, $x_{tW} = m_t^2/m_W^2$. The quantity Δ contains the dependence on the MSSM parameters and can be decomposed as

$$\Delta = \Delta_W + \Delta_H + \tilde{\Delta} \,. \tag{8}$$

In eq. (8), Δ_W denotes the Standard Model contribution, associated with the box diagrams involving the top quark and the W boson:

$$\Delta_W = A(x_{tW}) \,, \tag{9}$$

where the explicit expression of the function $A(x)$ is given in the appendix. Δ_H denotes the additional contributions from the box diagrams involving the physical charged Higgs boson of the MSSM [8]:

$$\Delta_H = \cot^4\beta\, x_{tH} \frac{1}{4} G(x_{tH}) + 2\cot^2\beta\, x_{tW} \left[F'(x_{tW}, x_{HW}) + \frac{1}{4} G'(x_{tW}, x_{HW}) \right] , \tag{10}$$

where $x_{tH} = m_t^2/m_{H^\pm}^2$, $x_{HW} = m_{H^\pm}^2/m_W^2$, $\tan\beta = v_2/v_1$, and the functions $G(x)$, $F'(x,y)$ and $G'(x,y)$ are given in the appendix. $\tilde{\Delta}$ denotes the contribution due to box diagrams

[2] In the MSSM there can be further corrections [5] to the relation between running and pole top quark mass, but we shall neglect them here.

with R-odd supersymmetric particles on the internal lines. Under our simplifying assumptions, we can take into account only the box diagram involving the $\tilde{t}_R$ and the charged higgsino. The general result of [9] then becomes

$$\tilde{\Delta} = \frac{x_{t\tilde{\chi}}}{4\sin^4\beta} G(x_{\tilde{t}\tilde{\chi}}, x_{\tilde{t}\tilde{\chi}}) \,, \tag{11}$$

where $x_{t\tilde{\chi}} = m_t^2/m_{\tilde{\chi}}^2$, $x_{\tilde{t}\tilde{\chi}} = m_{\tilde{t}}^2/m_{\tilde{\chi}}^2$, and $m_{\tilde{t}}$ ($m_{\tilde{\chi}}$) is the $\tilde{t}_R$ ($\tilde{H}^{\pm}$) mass.

Moving to the K^0–$\overline{K^0}$ system, the absolute value of the parameter ϵ_K is well approximated by the expression [7]:

$$|\epsilon_K| = \frac{2}{3} f_K^2 B_K \cdot \frac{m_K}{\sqrt{2}\Delta m_K} \cdot \left(\frac{\alpha_W}{4 m_W}\right)^2 \cdot x_{cW} \cdot |\Omega| \,, \tag{12}$$

where f_K is the K decay constant, B_K is the vacuum saturation parameter, Δm_K is the experimental K_L^0–K_S^0 mass difference. The quantity Ω, carrying the dependence on the mixing angles and the MSSM parameters, is given by [10]:

$$\Omega = \eta_{cc}\,\mathrm{Im}(K_{cs}K_{cd}^*)^2 + 2\eta_{ct}\,\mathrm{Im}(K_{cs}K_{cd}^*K_{ts}K_{td}^*)^2\,[B(x_{tW}) - \log x_{cW}] + \eta_{tt}\,\mathrm{Im}(K_{ts}K_{td}^*)^2\,x_{tc}\,\Delta \,, \tag{13}$$

where $\eta_{cc} \simeq 1.38$, $\eta_{ct} \simeq 0.47$ and $\eta_{tt} \simeq 0.57$ are QCD correction factors; $x_{cW} = m_c^2/m_W^2$, $x_{tc} = m_t^2/m_c^2$; the function $B(x)$ is given in the appendix; Δ is the same as in eq. (8), and contains all the dependence on the MSSM parameters. In principle, there are additional contributions due to charged Higgs exchange besides those appearing in Δ. However, for $\tan\beta \gtrsim 1$ they are much smaller than the standard contribution [10], hence they have been neglected[3].

We have studied the dependence of Δ on the parameters $(m_H, \tan\beta)$, characterizing the Higgs sector, and $(m_{\tilde{\chi}}, m_{\tilde{t}})$, characterizing the chargino-stop sector within our simplifying assumptions (similar studies were performed in [10, 11, 12]). It was already noticed in [10] that the interference between the three contributions in eq. (8) is always constructive, so that in general $\Delta_{MSSM} > \Delta_{SM}$. Besides the obvious symmetry due to the fact that $G(1/x) = xG(x)$, in the region of parameters of present phenomenological interest $\tilde{\Delta}$ is almost completely controlled by $m_{ave} \equiv (m_{\tilde{t}} + m_{\tilde{\chi}})/2$, with negligible dependence on $m_{\tilde{t}} - m_{\tilde{\chi}}$. Given the fact that in the MSSM, taking into account the present experimental bounds on the neutral Higgs bosons, $m_H \gtrsim 100$ GeV, for stops and charginos in the mass range accessible to LEP2, $\tilde{\Delta}$ dominates over Δ_H. Moreover, due to the additional enhancement factor $x_{t\tilde{\chi}}$, $\tilde{\Delta}$ represents the potentially largest contribution to Δm_{B_d}, and gives rise to a strong dependence on $\tan\beta$ near $\tan\beta = 1$, due to the $1/\sin^4\beta$ factor in eq. (11). Some quantitative information is given in fig. 2, which displays contours of the ratio

$$R_\Delta \equiv \frac{\Delta}{\Delta_W} \tag{14}$$

[3]In the fit to be described below, we have explicitly checked that the inclusion of such contributions does not modify the results appreciably.

in the plane $(\tan\beta, m_{ave})$, for $m_H = 100$ GeV (higher values of m_H do not displace significantly the contours, and we have taken for definiteness $m_{\tilde{t}} = m_{\tilde{\chi}}$). As can be seen, for values of $\tan\beta$ close to 1 and light stop and chargino, one can obtain $R_\Delta \gg 1$. However, a lower limit of $\tan\beta \gtrsim 1.5$ can be obtained by requiring that the top Yukawa coupling remain perturbative up to $M_{GUT} \sim 10^{16}$ GeV. One could also argue that charginos lighter than 65 GeV would have been copiously produced in the recent run of LEP 1.5, whilst no candidate events have been reported by the standard chargino searches [13]. However, the reader should keep in mind that no mass bound stronger than the LEP1 limit can be established yet if the chargino–neutralino mass difference is sufficiently small (a likely possibility in our approximations), or if charginos have R-parity violating decays with final states consisting of jets and no missing energy, or if the chargino production cross-section is suppressed by the destructive interference between the (γ, Z)-exchange and the $\tilde{\nu}_e$-exchange diagrams. For these reasons, we think that in our analysis we can safely consider chargino masses as low as 50 GeV or so. Thus, values of R_Δ as large as about 5 can still be obtained: we shall see in a moment how this compares with experimental data.

We now discuss the constraints coming from the measured values of Δm_{B_d} and ϵ_K. The dependence on the MSSM parameters is contained in the quantity Δ of eq. (8), so it would be desirable to obtain from the experimental data a bound on Δ. On the other hand, this requires some knowledge of the parameters characterizing the mixing matrix K. Notice that we cannot rely upon the SM fit to the matrix K, since among the experimental quantities entering this fit there are precisely Δm_{B_d} and ϵ_K, whose description now differs from the SM one.

We adopt here the Wolfenstein parametrization of the mixing matrix K:

$$K = \begin{pmatrix} 1 - \frac{\lambda^2}{2} & \lambda & A\lambda^3(\rho - i\eta) \\ -\lambda & 1 - \frac{\lambda^2}{2} & A\lambda^2 \\ A\lambda^3(1 - \rho - i\eta) & -A\lambda^2 & 1 \end{pmatrix} \tag{15}$$

The four experimental quantities used to constrain A, ρ and η are:

- The direct measure of the matrix element $|K_{cb}|$, from the semileptonic decay of the B meson [14]:
$$|K_{cb}| = 0.039 \pm 0.002\,. \tag{16}$$
This fixes the A parameter, and is not affected by the MSSM in any significant way.

- The direct measure of the ratio $|K_{ub}/K_{cb}|$ from the semileptonic charmless transitions of the B meson [7]:
$$|K_{ub}/K_{cb}| = 0.08 \pm 0.02\,. \tag{17}$$
This constrains the combination $\sqrt{\rho^2 + \eta^2}$, independently of the MSSM parameters.

- The B^0–$\overline{B^0}$ mass difference [15]:

$$\Delta m_{B_d} = (3.01 \pm 0.13) \times 10^{-13}\ \text{GeV}\,. \tag{18}$$

  This constrains the combination $A^2[(1-\rho)^2+\eta^2]$, as in the SM. However, it depends on the MSSM parameters through Δ.

- The parameter ϵ_K of CP violation [7]:

$$|\epsilon_K| = (2.26 \pm 0.02) \times 10^{-3}\,. \tag{19}$$

  Here one tests an independent combination of (A,ρ,η), which depends on the value of Δ in the MSSM.

To derive the desired bound on Δ, we have performed a fit to these data suitable for the MSSM, i.e. keeping A, ρ, η and the Δ as independent variables. The results of the fit are sensitive to the input values of the parameters $f^2_{B_d}B_{B_d}$ and B_K. We have repeated the fit for various allowed values of $f^2_{B_d}B_{B_d}$ and B_K, to estimate the effect of the corresponding theoretical uncertainties. We have checked that, by fixing Δ to its SM value, $\Delta = \Delta_W = 0.551$ for $m_t = 170$ GeV, we recover the results for (A,ρ,η) of the SM fit [7].

$\sqrt{f^2_{B_d}B_{B_d}}$ (MeV)	ρ	η	Δ
160	-0.19 ± 0.14	0.31 ± 0.04	0.55 ± 0.15
160	$+0.30^{+0.12}_{-0.61}$	$0.21^{+0.14}_{-0.04}$	$1.52^{+0.70}_{-1.10}$
180	-0.08 ± 0.23	0.35 ± 0.05	0.50 ± 0.22
180	$+0.23 \pm 0.22$	0.28 ± 0.08	0.97 ± 0.52
200	$+0.09 \pm 0.41$	0.37 ± 0.11	0.54 ± 0.46
220	$+0.10^{+0.21}_{-0.22}$	$0.41^{+0.06}_{-0.07}$	$0.43^{+0.27}_{-0.15}$
240	$+0.10^{+0.18}_{-0.19}$	$0.45^{+0.06}_{-0.07}$	$0.34^{+0.17}_{-0.10}$

Table 1: Results of the fit for $m_t = 170$ GeV, $B_K = 0.75$ and for different values of $f^2_{B_d}B_{B_d}$. The fitted value of A ranges from 0.80 to 0.83, with an uncertainty of 0.04.

In table 1 we show our results for several choices of $f^2_{B_d}B_{B_d}$ and for $B_K = 0.75$. In all cases the parameter A, basically determined by $|K_{cb}|$, essentially coincides with its SM determination. For relatively low values of $f^2_{B_d}B_{B_d}$, small values of η are preferred and the χ^2 function has two minima: this is due to the constraint coming from $|K_{ub}/K_{cb}|$, which is sensitive to $\sqrt{\rho^2+\eta^2}$ and, for positive η, has a twofold ambiguity in ρ. For $\sqrt{f^2_{B_d}B_{B_d}} \lesssim 190$ MeV, the negative ρ solution is the one favoured by the SM and leads to a central value for Δ which is very close to the SM one. On the contrary, the positive ρ solution, which is absent in the SM for the current choice of parameters, corresponds to

a higher Δ. For relatively high values of $f_{B_d}^2 B_{B_d}$, the χ^2 function has a unique minimum, ρ is very close to zero and Δ is close to its SM value. It is clear that, qualitatively, large values of Δ can be allowed only for small $f_{B_d}^2 B_{B_d}$.

In table 2 we show the influence of the B_K parameter. The highest values of Δ are obtained when $B_K = 0.9$. This can be qualitatively understood as follows: a large Δ is compatible with the measured Δm_{B_d} only when η is close to zero. On the other hand, since $|\epsilon_K|$ is proportional to η, the smallness of η should be compensated by B_K, which is then required to be large.

$\sqrt{f_{B_d}^2 B_{B_d}}$ (MeV)	160	180	200	220	240
$B_K = 0.6$	$1.12^{+0.73}_{-0.65}$	0.67 ± 0.39	$0.51^{+0.27}_{-0.16}$	$0.39^{+0.17}_{-0.11}$	
$B_K = 0.9$	$1.69^{+0.72}_{-0.50}$	$1.20^{+0.54}_{-0.86}$	$0.82^{+0.43}_{-0.39}$	0.46 ± 1.54	$0.36^{+0.26}_{-0.13}$

Table 2: Central values and 1σ errors for Δ, for different choices of $f_{B_d}^2 B_{B_d}$ and B_K. When two minima are present, only the largest central value for Δ is quoted.

It is not straightforward to translate the above results into a single definite bound on Δ, or, equivalently, on $R_\Delta = \Delta/\Delta_W \simeq 1.8\Delta$. Values of R_Δ as large as 5 (see fig. 2) are clearly disfavoured, but cannot be rigorously excluded if one keeps in mind the theoretical uncertainties on the parameters $f_{B_d}^2 B_{B_d}$ and B_K. For the moment, we can only state that small $\tan\beta$ and very light chargino and stop require small $f_{B_d}^2 B_{B_d}$ and large B_K, and imply small η and positive and large ρ.

The expression of Δm_{B_s} can be trivially obtained from eq. (1) by making everywhere the replacement $d \to s$. The present 95% CL limit [15], $\Delta m_{B_s} \geq 4.0 \times 10^{-12}$ GeV, does not provide additional constraints on Δ. One obtains:

$$x_{tW}\Delta \geq 1.14 \left(\frac{230}{f_{B_s}\sqrt{B_{B_s}}\ (\mathrm{MeV})} \right)^2 \left(\frac{0.8}{A} \right)^2 . \tag{20}$$

In the SM, for $m_t = 170$ GeV one has $x_{tW}\Delta_W = 2.47$. Since $\Delta \geq \Delta_W$, the previous limit is always respected in the MSSM, for all values of the parameters. On the other hand, we can derive some information on Δm_{B_s} in the MSSM from the relation:

$$\frac{\Delta m_{B_s}}{\Delta m_{B_d}} = \frac{m_{B_s}}{m_{B_d}} \cdot \xi_s^2 \cdot \frac{\eta_{B_s}}{\eta_{B_d}} \left| \frac{K_{ts}}{K_{td}} \right|^2 , \tag{21}$$

where $\xi_s = [f_{B_s}\sqrt{B_{B_s}}]/[f_{B_d}\sqrt{B_{B_d}}]$. One expects $\xi_s = 1.16 \pm 0.1$ [16] and $\eta_{B_s} = \eta_{B_d}$. Then, from eq. (21) one obtains:

$$\Delta m_{B_s} \simeq (28 \pm 5) \cdot \frac{\Delta m_{B_d}}{[(1-\rho)^2 + \eta^2]} . \tag{22}$$

This relation is valid both in the SM and in the limit of the MSSM considered here. However, the high value of Δ which could be obtained in the MSSM for small $\tan\beta$ and

light chargino and stop, would imply a value for the combination $(1-\rho)^2+\eta^2$ smaller than in the SM, as can be seen from table 1. On this basis, we conclude that the value expected for Δm_{B_s} in the MSSM, when $\tan\beta$ is small and stop and chargino are both light, is always higher than the one foreseen in the SM. However, in view of the existing uncertainties on ρ and η, a more precise estimate of Δm_{B_s} in the MSSM is not yet possible.

To conclude this section, we would like to comment on the MSSM effects on the ratio ϵ'/ϵ. These have been analysed, at leading order in QCD and QED, in ref. [17]. We recall that, on the experimental side, there are two independent results for Re (ϵ'/ϵ):

$$\begin{array}{rcll} \mathrm{Re}\,(\epsilon'/\epsilon) & = & 23\pm 6.5\times 10^{-4} & \mathrm{NA31}\quad [18]\,, \\ \mathrm{Re}\,(\epsilon'/\epsilon) & = & 7.4\pm 6.0\times 10^{-4} & \mathrm{E731}\quad [19]\,. \end{array} \tag{23}$$

The SM value of $\mathrm{Re}\,\epsilon'/\epsilon$ is typically of order 10^{-4} for $m_t = 150$–190 GeV, decreases for increasing m_t, and vanishes for $m_t = 200$–220 GeV. In the MSSM, it is possible to enhance the SM prediction by at most 40% for $m_t \simeq 170$ GeV and up to 60% for $m_t \simeq 190$ GeV. The enhancement is attained for chargino and stop masses close to the present LEP limit, with the other squarks and the charged Higgs much heavier. This modest enhancement cannot explain the large central value of $\mathrm{Re}\,\epsilon'/\epsilon$ suggested by the NA31 experiment and, on the other hand, is perfectly compatible with the data of the E731 collaboration. A reduction of $\mathrm{Re}\,\epsilon'/\epsilon$ with respect to the SM value is also achievable in the MSSM. This requires a light charged Higgs and light charginos and stops. Part of the effect is due to the fact that $\mathrm{Re}\,\epsilon'/\epsilon$ is proportional to η, which, as discussed above, can be considerably smaller than in the SM. In this case a vanishing or even negative value of $\mathrm{Re}\,\epsilon'/\epsilon$ can be obtained for $m_t = 150$–190 GeV. This depletion, which potentially represents the most conspicuous effect of minimal supersymmetry, is however very difficult to test, due to the insufficient experimental sensitivity. In conclusion, the present data on ϵ'/ϵ do not provide any additional constraint on the MSSM parameter space.

$\boxed{b \to s\gamma}$

The recent CLEO result [20] on the inclusive $B \to X_s\gamma$ decay, $BR(B \to X_s\gamma) = (2.32 \pm 0.67)\times 10^{-4}$, agrees with the SM predictions based on the partonic process $b \to s\gamma$ (for a review and references, see e.g. [21]), and at the same time constrains possible extensions of the SM, in particular the MSSM. A very conservative estimate [22] gives $BR(B \to X_s\gamma)_{SM} = (2.55 \pm 1.28)\times 10^{-4}$, whereas other authors [23] quote similar central values but smaller errors, at the level of 30%. Under our simplifying assumptions, the $b \to s\gamma$ amplitude at a scale $\mathcal{O}(M_W)$ receives additional contributions from top and charged Higgs (stop and chargino) exchange, which interfere constructively (destructively) with the SM contributions, dominated by top and W exchange. The amplitude at a scale $\mathcal{O}(m_b)$ gets both multiplicatively and additively renormalized by QCD corrections. The latter effect is mainly due to the mixing between the magnetic operator (O_7) and a four-quark operator

(O_2). We will express our results in terms of the ratio

$$R_\gamma \equiv \frac{Br(B \to X_s\gamma)_{MSSM}}{Br(B \to X_s\gamma)_{SM}}\,, \tag{24}$$

which we identify with the corresponding ratio of $b \to s\gamma$ squared amplitudes. Then we estimate:

$$R_\gamma \simeq \left[\frac{C(A_W + A_H + \tilde{A}) + D}{C\ A_W + D}\right]^2 , \tag{25}$$

where [9] $C \simeq 0.66$, $D \simeq 0.35$,

$$A_W = x_{tW}(2F_1(x_{tW}) + 3F_2(x_{tW}))\,, \tag{26}$$

$$A_H = x_{tH}\left\{\cot^2\beta\left[\frac{2}{3}F_1(x_{tH}) + F_2(x_{tH})\right] + \left[\frac{2}{3}F_3(x_{tH}) + F_4(x_{tH})\right]\right\} , \tag{27}$$

$$\tilde{A} = -\frac{x_{t\tilde{t}}}{\sin^2\beta}\left[F_1(x_{\tilde{\chi}\tilde{t}}) + \frac{2}{3}F_2(x_{\tilde{\chi}\tilde{t}})\right] . \tag{28}$$

Similarly to the previously discussed R_Δ, we have found that R_γ depends on $m_{\tilde{t}}$ and $m_{\tilde{\chi}}$ essentially through their sum (not their difference). Then we can focus as before on the variable $m_{ave} = (m_{\tilde{t}}+m_{\tilde{\chi}})/2$. Figs. 3 and 4 show contour lines of R_γ, in the (m_{ave}, m_H) plane for $\tan\beta = 1.5, 5$ and in the $(m_{ave}, \tan\beta)$ plane for $m_H = 100, 500$ GeV, respectively, taking for definiteness $m_{\tilde{t}} = m_{\tilde{\chi}}$. The contour $R_\gamma = 1$ corresponds to the situations in which the 'extra' contributions A_H and $\tilde{A}$ cancel against each other (the possibility of these cancellations was emphasized in ref. [24]), so that the SM result is recovered. Since the comparison between theory and experiment is dominated by the theoretical error, the allowed region can be estimated conservatively to be $0.5 \lesssim R_\gamma \lesssim 1.5$, or slightly less conservatively $0.7 \lesssim R_\gamma \lesssim 1.3$. In fig. 3, one can notice the strong positive correlation between m_{ave} and m_H. In other words, light charged Higgs and heavy stop and chargino would give too large a value for $BR(B \to X_s\gamma)$, whereas heavy charged Higgs and light stop and chargino would give too small a value. In addition, fig. 4 shows a moderate dependence on $\tan\beta$ in the range $1 \lesssim \tan\beta \lesssim 2$.

$\boxed{R_b}$

The calculation of $R_b \equiv \Gamma(Z \to b\bar{b})/\Gamma(Z \to \text{hadrons})$ in the MSSM was performed in [25]. Specializing those results to the limiting case under discussion, we can write

$$R_b = (R_b)_{SM}\left[1 + 0.78 \times \frac{\alpha_W}{2\pi}\frac{v_L}{v_L^2 + v_R^2}(F_H + \tilde{F})\right] , \tag{29}$$

where, for the input values $M_t = 180 \pm 12$ GeV and $\alpha_S(m_Z) = 0.125 \pm 0.007$,

$$(R_b)_{SM} = 0.2156 \pm 0.0005\,, \tag{30}$$

and

$$v_L = -\frac{1}{2} + \frac{1}{3}\sin^2\theta_W\,, \qquad v_R = \frac{1}{3}\sin^2\theta_W\,. \tag{31}$$

F_H and $\tilde{F}$ are associated with top-Higgs and stop-higgsino loops, respectively, and read

$$\begin{aligned} F_H = \Big\{ & b_1(m_H, m_t) v_L - c_0(m_t, m_H) v_L^{(H)} + m_t^2 c_2(m_H, m_t) v_L^{(t)} \\ & + \Big[m_Z^2 c_6(m_H, m_t) - \tfrac{1}{2} - c_0(m_H, m_t) \Big] v_R^{(t)} \Big\} \lambda_H^2\,, \end{aligned} \tag{32}$$

$$\begin{aligned} \tilde{F} = \Big\{ & b_1(m_{\tilde{t}}, m_{\tilde{\chi}}) v_L - c_0(m_{\tilde{\chi}}, m_{\tilde{t}}) v_R^{(t)} \\ + \Big[& m_Z^2 c_6(m_{\tilde{t}}, m_{\tilde{\chi}}) - \tfrac{1}{2} - c_0(m_{\tilde{t}}, m_{\tilde{\chi}}) + m_{\tilde{\chi}}{}^2 c_2(m_{\tilde{t}}, m_{\tilde{\chi}}) \Big] v_L^{(H)} \Big\} \tilde{\lambda}^2\,, \end{aligned} \tag{33}$$

where the functions b_1, c_0, c_2 and c_6 are given in the appendix, and

$$\lambda_H = \frac{m_t}{\sqrt{2} m_W \tan\beta}\,, \qquad \tilde{\lambda} = \frac{m_t}{\sqrt{2} m_W \sin\beta}\,, \tag{34}$$

$$v_L^{(t)} = \frac{1}{2} - \frac{2}{3}\sin^2\theta_W\,, \qquad v_R^{(t)} = -\frac{2}{3}\sin^2\theta_W\,, \qquad v_L^{(H)} = -\frac{1}{2} + \sin^2\theta_W\,. \tag{35}$$

A quantitative estimate of the possible effects is given in fig. 5, which shows contours of R_b in the $(m_{\tilde{\chi}}, m_{\tilde{t}})$ plane, for some representative values of $\tan\beta$ and m_H. One can see that, in our limiting case, values of R_b as high as 0.218 can be reached, for $\tilde{t}_R$ and higgsinos very close to 50 GeV. Notice also that the dependence of R_b on $m_{\tilde{\chi}}$ is stronger than the dependence on $m_{\tilde{t}}$, which makes the higgsino mass μ the most relevant parameter. The dependences on $\tan\beta$ and on m_H are not very strong, and the effect is maximal for $m_{\tilde{\chi}}$, $m_{\tilde{t}}$ as close as possible to their experimental limits, low $\tan\beta$ (maximal top Yukawa coupling) and high m_H (minimal negative interference with the charged Higgs loops).

In the past, it was suggested [2] that an improved fit to $\alpha_S(m_Z)$ and to R_b could be obtained by allowing for some new physics that enhances R_b with respect to its SM prediction. The most recent experimental data [26] give $R_b = 0.2219 \pm 0.0017$, with a strong positive correlation with $R_c = 0.1540 \pm 0.0074$, which also significantly deviates from its SM prediction, $(R_c)_{SM} = 0.1724 \pm 0.0003$. If one fixes R_c to its SM value, the fit to the LEP data gives $R_b = 0.2205 \pm 0.0016$. Even in the last, most favourable case, our limiting case of the MSSM cannot produce R_b closer than 1.5σ to its experimental value. Slightly better agreement can be obtained for very large values of $\tan\beta$ and A^0 as light as allowed by the present experimental limits, but this case cannot be quantitatively studied within the present approximations.

$$\boxed{t \to \tilde{t}\tilde{H}_S,\ \tilde{t}\tilde{H}_A,\ bH^+}$$

In the presence of sufficiently light charged Higgs boson, stop and higgsinos, new decay modes are kinematically accessible in the top quark decays, in addition to the standard

mode $t \to bW^+$: assuming heavy sbottom squarks, they are $t \to \tilde{t}\tilde{H}_S$, $\tilde{t}\tilde{H}_A$, bH^+. The corresponding partial widths are reported below [27]:

$$\Gamma(t \to \tilde{t}\tilde{H}_i^0) = \frac{\sqrt{[m_t^2 - (m_{\tilde{t}} + m_{\tilde{H}_i})^2][m_t^2 - (m_{\tilde{t}} - m_{\tilde{H}_i})^2]}}{16\pi m_t^3} \cdot \mathcal{A}_i \,, \tag{36}$$

$$\mathcal{A}_i = \frac{g^2 m_t^2}{8 m_W^2 \sin^2\beta}(m_t^2 + m_{\tilde{H}_i}^2 - m_{\tilde{t}}^2) \,, \qquad (i = S, A) \,; \tag{37}$$

$$\Gamma(t \to bH^+) = \frac{\sqrt{[m_t^2 - (m_H + m_b)^2][m_t^2 - (m_H - m_b)^2]}}{16\pi m_t^3} \cdot \mathcal{A}_H \,, \tag{38}$$

$$\mathcal{A}_H = \frac{g^2 m_t^2}{4 m_W^2 \tan^2\beta}(m_t^2 + m_b^2 - m_H^2) \,; \tag{39}$$

$$\Gamma(t \to bW^+) = \frac{\sqrt{[m_t^2 - (m_W + m_b)^2][m_t^2 - (m_W - m_b)^2]}}{16\pi m_t^3} \cdot \mathcal{A}_W \,, \tag{40}$$

$$\mathcal{A}_W = \frac{g^2}{4}\left[m_t^2 + m_b^2 - m_W^2 - \frac{m_W^4 - (m_t^2 - m_b^2)^2}{m_W^2}\right] \,. \tag{41}$$

With the help of fig. 6, which displays contours of $BR(t \to bW^+)$ in the $(\mu, m_{\tilde{t}})$ plane, for some representative values of $\tan\beta$ and m_H, we can see that deviations from the SM prediction $BR(t \to bW^+) \simeq 1$ can be very significant, up to $BR(t \to bW^+) \sim 0.4$. However, this cannot be transformed easily into a constraint on the parameter space, as attempted in [28]: first, the perturbations to our limiting case, illustrated in fig. 1, can modify the results for the top branching ratios, but cannot be accounted for without introducing additional parameters such as M_2; second, it is not clear how strong a lower bound the present CDF and D0 data can provide on $BR(t \to bW^+)$: deriving such a bound requires not only the detailed knowledge of the experimental selection criteria, but also assumptions about the production cross-section and the stop and higgsinos branching ratios. We do not feel in a position to do so reliably, so we content ourselves with displaying the contours in fig. 6. As a reference value for the CDF and D0 sensitivity, we can tentatively take $BR(t \to bW^+) = 0.7$: it is then clear that the foreseeable Tevatron bounds on exotic top decays will significantly constrain the light $\tilde{t}_R$-higgsino scenario, in qualitative agreement with the conclusions of ref. [28].

3. Higgsinos and $\tilde{t}_R$ are among the most plausible light supersymmetric particles, and a mass range for these states within the experimental reach of LEP2 would represent an appealing scenario. At the moment, this range is still compatible with the precision electroweak data, and it could even play a role in reducing the present discrepancy concerning R_b. A more systematic analysis is however required to confront this possibility with the available experimental information, including the rich input coming from flavour-changing transitions. The aim of this note has been to improve the existing studies, which often

focus on a single specific process, by discussing all the relevant constraints, albeit in a simplified setting.

Among the observables that are potentially most sensitive, we focused on the mixing parameter Δm_{B_d} and the CP–violating parameter ϵ_K. One could have expected that, given the present experimental precision on those data, light $\tilde{t}_R$ and higgsinos could already be ruled out, at least for small $\tan\beta$. Actually, due to the theoretical uncertainties affecting $f^2_{B_d} B_{B_d}$ and B_K, we have been led to a milder statement. The corner in the MSSM parameter space with very light $\tilde{t}_R$ and higgsino, and small $\tan\beta$, requires small values of $f^2_{B_d} B_{B_d}$ and large ones for B_K, at the border of the presently allowed theoretical ranges. Moreover, positive and large values of ρ and small values of η are preferred, with immediate implications on the values of the CP–violating asymmetries in B decays at future facilities.

The process that gives the most significant constraints on the reduced MSSM parameter space, corresponding to the limiting case of light $\tilde{t}_R$ and higgsinos, is $b \to s\gamma$ (see figs. 3 and 4). For example, the existence of $\tilde{t}_R$ and higgsinos around 60 GeV would require[4] $m_H \lesssim 100$ GeV for $\tan\beta = 1.5$, $m_H \lesssim 200$ GeV for $\tan\beta = 5$. This has also indirect effects on the allowed values for the neutral Higgs bosons of the MSSM. Since in our limiting case the sum rule $m^2_A + m^2_W = m^2_H$ remains valid to quite a good accuracy after the inclusion of radiative corrections [29], one gets corresponding approximate upper bounds on m_A. The mass of the lightest CP-even state could also be affected, since its tree-level value depends on $(m_A, \tan\beta)$, whilst radiative corrections [30] are mainly controlled by the logarithmic dependence on $m_{\tilde{t}_1} m_{\tilde{t}_2}$, for fixed M_t and $\tan\beta$. Under our assumptions, however, we are still free to push $m_{\tilde{t}_L}$ to values sufficiently high that the experimental bounds can be evaded.

The measurement of R_b at LEP and the study of top decays at the Tevatron cannot be transformed, for the moment, into precise bounds on the MSSM parameter space. In the case of R_b, besides the open question of the correlation with R_c, the size of the typical effects of light $\tilde{t}_R$ and higgsinos is considerably smaller than the discrepancy between the SM prediction and the experimental average. In the case of top decays, only the CDF and D0 collaborations have the appropriate tools to establish reliable bounds on exotic channels, either directly or by the extraction of $BR(t \to bW^+)$. If, as expected, the bound setlles around $BR(t \to bW^+) > 0.7$, then the surviving region of the $(m_{\tilde{\chi}}, m_{\tilde{t}})$ plane will allow at most for $\Delta R_b \lesssim 10^{-3}$, a rather marginal improvement over the SM when compared with the experimental data.

Finally, we recall that the parameter space discussed in section 2 starts to get significant constraints from the direct searches for stops and charginos, both at the Tevatron [31] and at LEP [13, 32]. We hope that the analysis reported in this paper will contribute to the understanding of the interplay between indirect and direct signals of light $\tilde{t}_R$ and higgsinos.

[4]These bounds could be somewhat relaxed, however, when mixing effects in the chargino and stop sectors are included.

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Appendix

We collect in this appendix the functions, obtained from one-loop diagrams, appearing in the formulae given in the text.

$$A(x) = \frac{1}{4(x-1)^3}(x^3 - 12x^2 + 15x - 4 + 6x^2 \log x)\,, \tag{42}$$

$$G(x) = \frac{-1 + x^2 - 2x \log x}{(x-1)^3}\,, \tag{43}$$

$$F'(x,y) = \frac{-1 + x - \log x}{(x-1)^2(y-x)} + \frac{\dfrac{x \log x}{x-1} + \dfrac{y \log y}{1-y}}{(x-y)^2}\,, \tag{44}$$

$$\begin{aligned} G'(x,y) &= \frac{3 - 4x + x^2 + 4x \log x - 2x^2 \log x}{2(x-1)^2(y-x)} \\ &- \frac{\dfrac{3(y-x)}{2} + \dfrac{x^2 \log x}{x-1} + \dfrac{y^2 \log y}{1-y}}{(x-y)^2}\,, \end{aligned} \tag{45}$$

$$B(x) = \log x - \frac{3}{4}\frac{x}{x-1}\left(-1 + \frac{x}{x-1}\log x\right)\,, \tag{46}$$

$$F_1(x) = \frac{1}{12(x-1)^4}(x^3 - 6x^2 + 3x + 2 + 6x \log x)\,, \tag{47}$$

$$F_2(x) = \frac{1}{12(x-1)^4}(2x^3 + 3x^2 - 6x + 1 - 6x^2 \log x)\,, \tag{48}$$

$$F_3(x) = \frac{1}{2(x-1)^3}(x^2 - 4x + 3 + 2 \log x)\,, \tag{49}$$

$$F_4(x) = \frac{1}{2(x-1)^3}(x^2 - 1 - 2x \log x)\,. \tag{50}$$

$$b_1(m_1, m_2) = -\frac{1}{4} + \frac{m_2^2}{2(m_1^2 - m_2^2)} + \frac{\log(m_1^2/\mu^2)}{2} + \frac{m_2^4 \log(m_2^2/m_1^2)}{2(m_1^2 - m_2^2)^2}\,, \tag{51}$$

$$\begin{aligned} c_0(m_1, m_2) = &\int_0^1 dx \left\{ -\frac{(m_1^2 - m_1^2 x + m_2^2 x)[1 - \log(m_1^2 - m_1^2 x + m_2^2 x)/\mu^2]}{m_1^2 - m_2^2 + m_Z^2 x} \right. \\ &\left. + \frac{(m_2^2 - m_Z^2 x + m_Z^2 x^2)[1 - \log(m_2^2 - m_Z^2 x + m_Z^2 x^2)/\mu^2]}{m_1^2 - m_2^2 + m_Z^2 x} \right\}\,, \end{aligned} \tag{52}$$

$$c_2(m_1, m_2) = \int_0^1 dx \frac{\log[(m_1^2 - m_1^2 x + m_2^2 x)/(m_2^2 - m_Z^2 x + m_Z^2 x^2)]}{m_1^2 - m_2^2 + m_Z^2 x}\,, \tag{53}$$

$$c_6(m_1, m_2) = \int_0^1 dx \frac{x \log[(m_1^2 - m_1^2 x + m_2^2 x)/(m_2^2 - m_Z^2 x + m_Z^2 x^2)]}{m_1^2 - m_2^2 + m_Z^2 x}\,. \tag{54}$$

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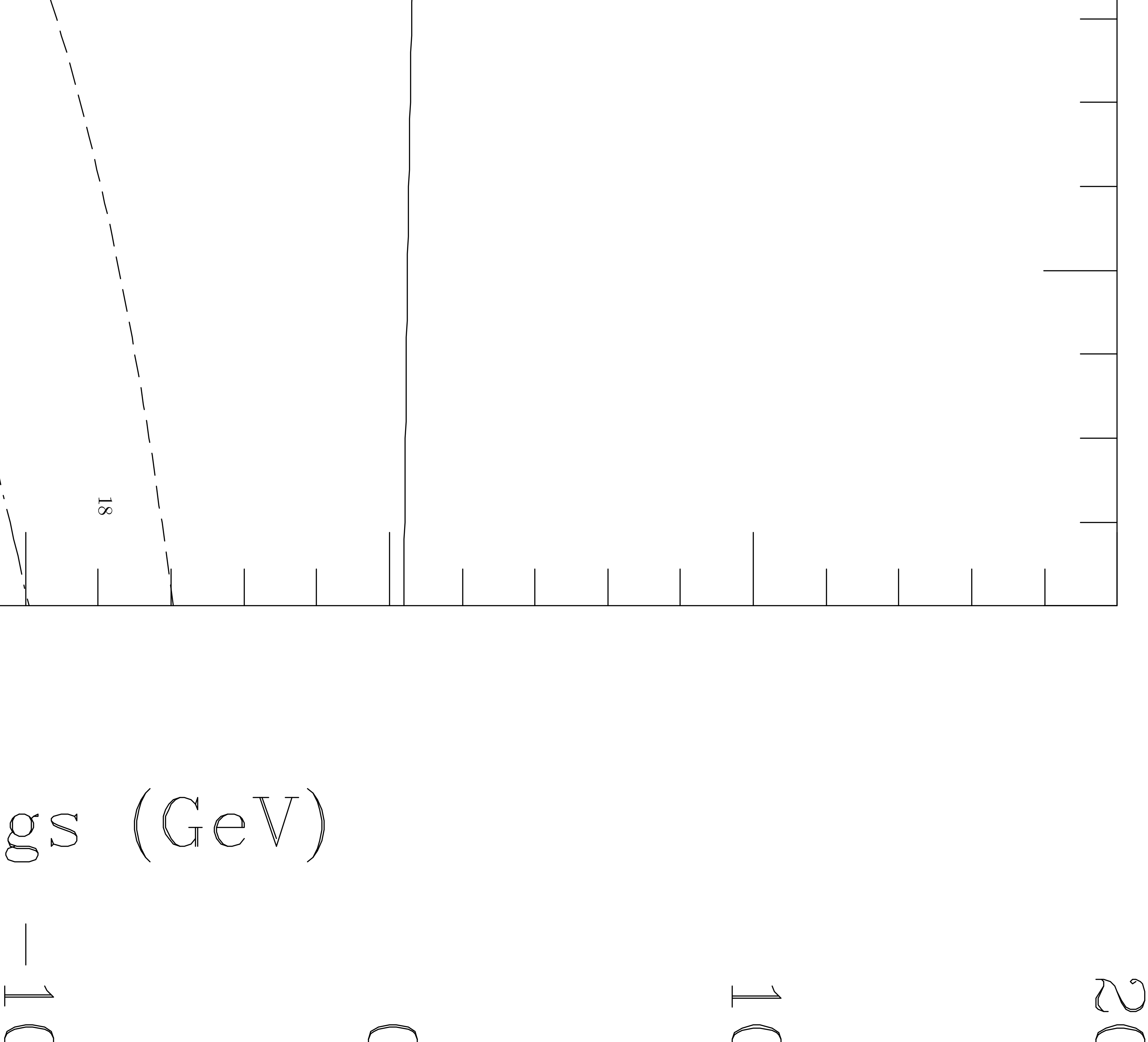
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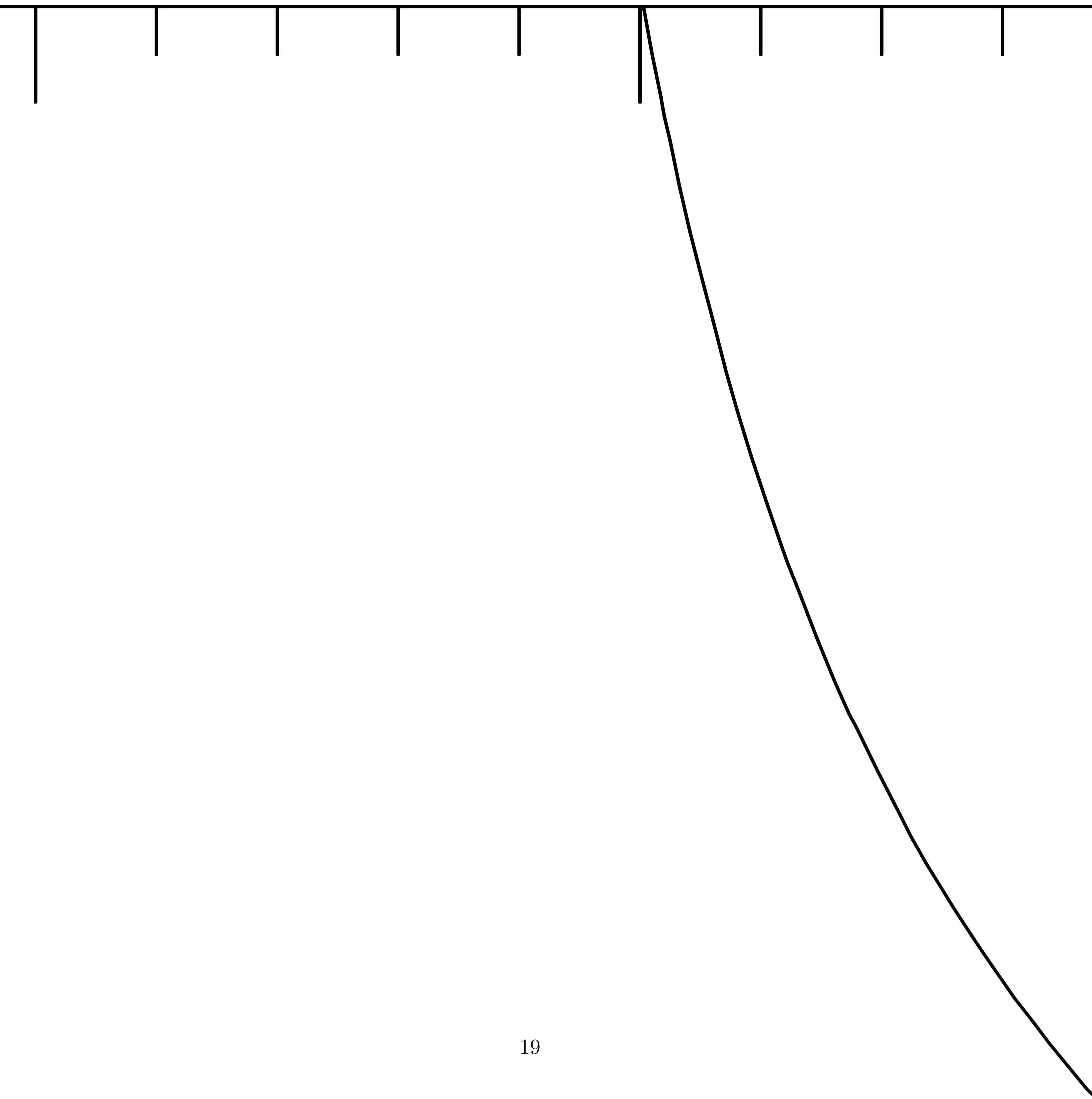

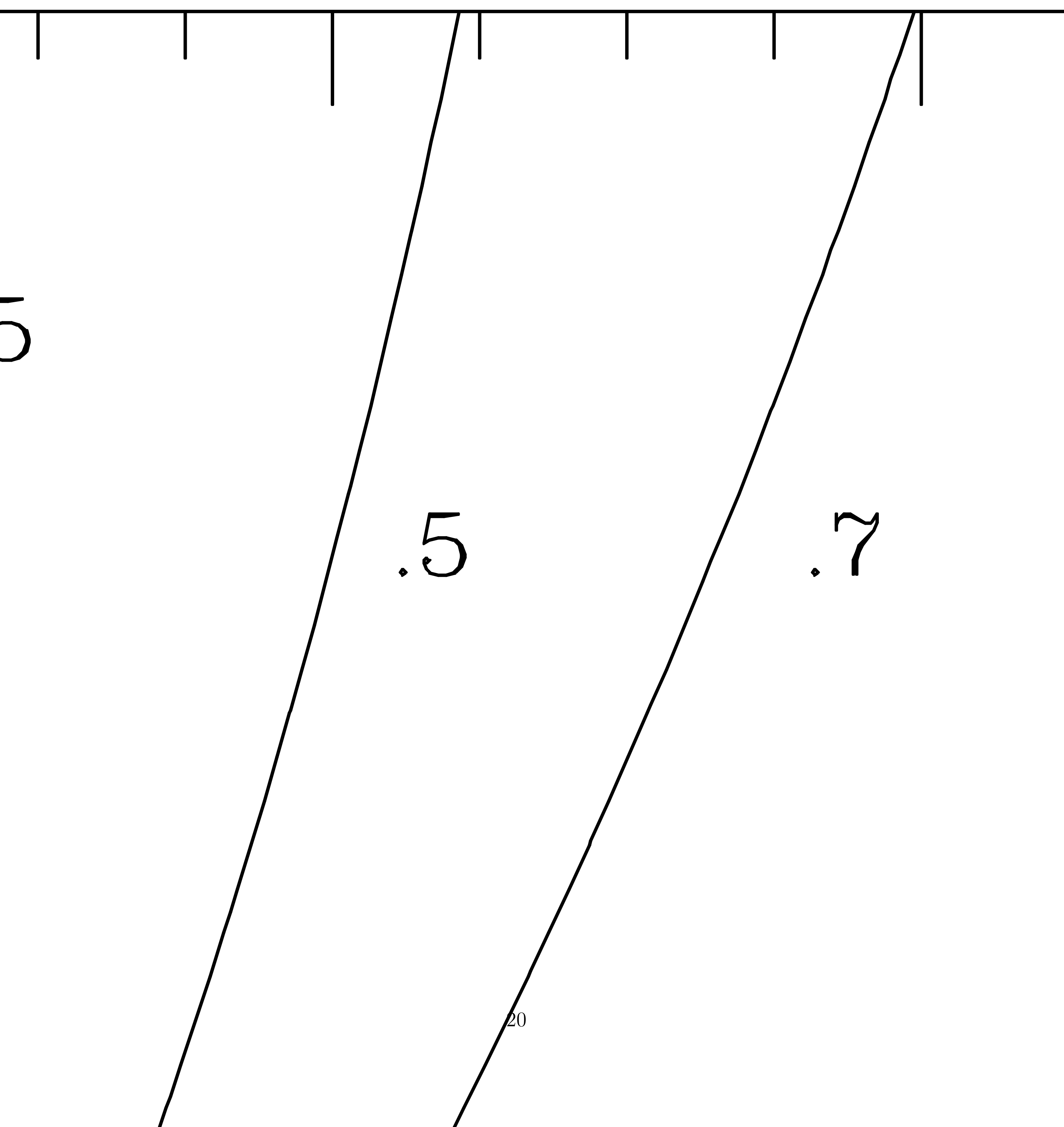
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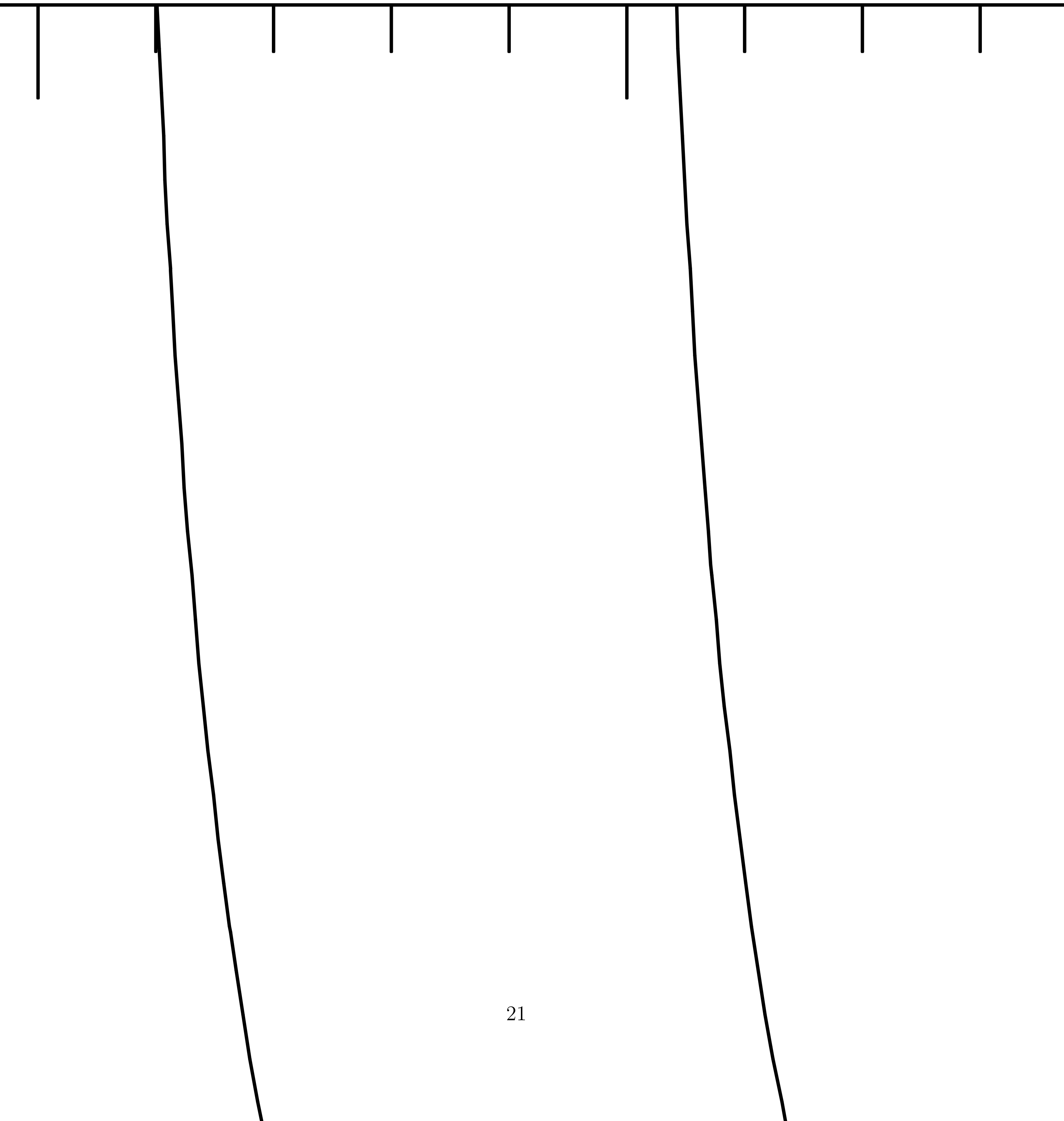

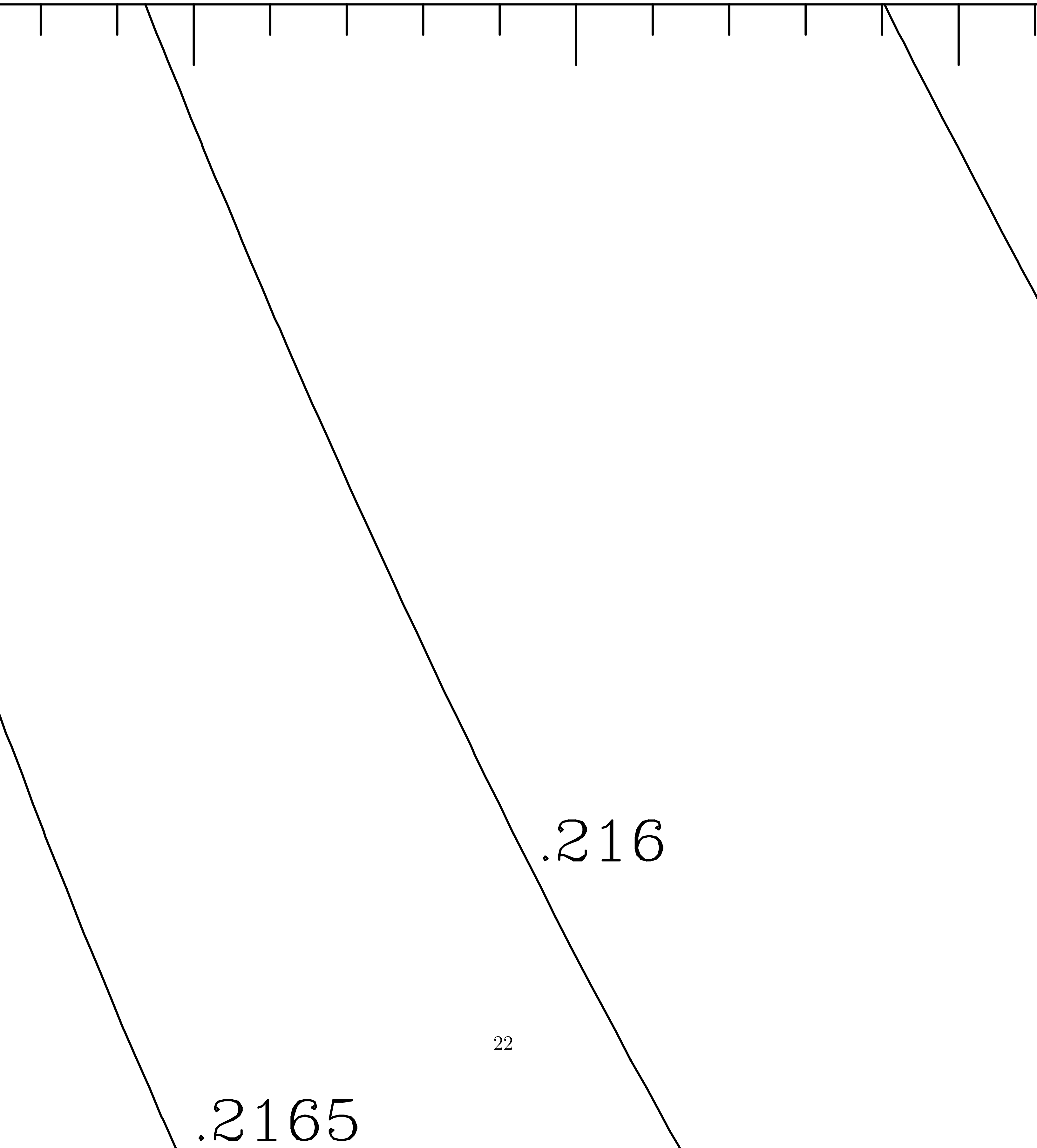
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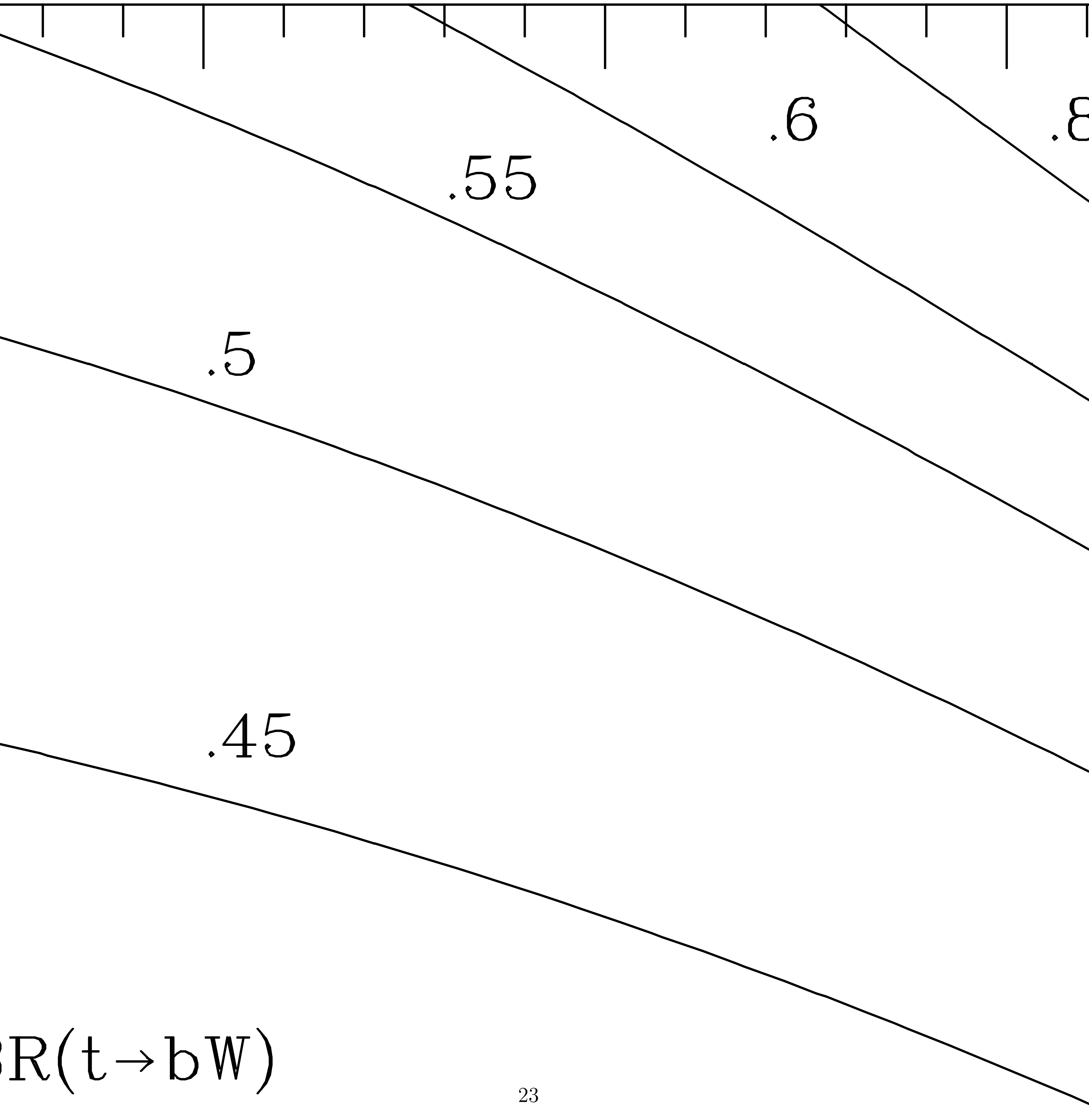
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